

Majorana versus Dirac Mass from Holomorphic Supersymmetric Nambu-Jona-Lasinio Model

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Abstract

We study the theoretical features in relation to dynamical mass generation and symmetry breaking for the recently proposed holomorphic supersymmetric Nambu–Jona-Lasinio model. The basic model has two different chiral superfields with a strongly coupled dimension five four-superfield interaction. In addition to the possibility of generation of Dirac mass between the pair established earlier, we show here the new option of generation of Majorana masses for each chiral superfield. We also give a first look at what condition may prefer Dirac over Majorana mass, illustrating that a split in the soft supersymmetry breaking masses is crucial. In particular, in the limit one of the soft mass vanishes, we show that generation of the Majorana mass is no longer an option, while the Dirac mass generation survives well. The latter is sensitive mostly to the average of the two soft masses. The result has positive implication on the application of the model framework towards dynamical electroweak symmetry breaking with Higgs superfields as composites.

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I. INTRODUCTION

Dynamical mass generation and symmetry breaking is a very interesting theoretical topic with important phenomenological applications. One of the simplest model of the kind is the Nambu–Jona-Lasinio (NJL) model [1]. It is also the first explicit model of spontaneous symmetry breaking. Analysis of the nonperturbative gap equation established that with a strong enough four-fermion interaction, a symmetry breaking Dirac fermion mass would be resulted. When applied to the electroweak symmetry breaking of the Standard Model, it fails to give the relatively small experimental top quark mass [2]. Introducing heavier fourth family quarks to take the role of the top is essentially ruled out by other experimental constraints. The more interesting option of supplementing supersymmetry [3, 4] requires too low a $\tan\beta$ value to stay phenomenologically viable. The latter situation can be resolved with an alternative supersymmetrization of the NJL model recently proposed [5]. The version has a dimension five four-superfield interaction, which otherwise mimics well most basic features of the NJL model.

Supersymmetry is an important theme in modern physics. One of specially attractive feature, in our opinion, is that scalar fields are now part of the chiral superfields with the chiral fermions. The chirality forbids any gauge invariant mass before breaking any symmetry. Moreover, the full matter (super)field spectrum is now strongly constrained by the gauge symmetry and their anomaly cancellation conditions. Introduction of the vectorlike pair of Higgs superfields with their un-natural gauge invariant mass in the usual formulation of the supersymmetric Standard Model looks particularly unattractive from the theoretical perspective. An NJL mechanism, with the Higgs superfield(s) generated as composite and the electroweak scale generated by strong dynamics is hence very appealing. The Holomorphic Supersymmetric Nambu–Jona-Lasinio model (HSNJL) construction [5, 6] gives exactly such a scenario that looks compatible with all known experimental constraints. In Ref.[6], superfield gap equation analyses of the Dirac mass generation have been performed for the HSNJL model and the old supersymmetric model. Nontrivial symmetry breaking masses were established.

Distinguished from the old models, the HSNJL model is very rich in interesting theoretical features, some of which we report here. Firstly, the HSNJL model is capable of generating also Majorana masses of the chiral superfields. Note that the basic model has two superfields,

which could otherwise be the Dirac pair. In fact, in the generic case, the story of dynamical mass generation and hence resulting symmetry breaking pattern becomes more complicated than the naive Dirac mass generation analysis would otherwise conclude. We present in this letter an illustration of the Majorana mass generation – a feature that is unknown for models in the literature. Short of doing a comprehensive and fully generic gap equation analysis, we will compare the Majorana mass result here versus the Dirac mass result in our previous paper [6] to give a first answer to the competition of Majorana versus Dirac superfield masses. We will discuss how a splitting between the two soft supersymmetry breaking masses favors the generation of Dirac superfield mass. In a very interesting particular case, we will show that at the limit one of the soft masses vanishes, the nontrivial Majorana mass would be killed as nontrivial Dirac mass solution survives. Dynamical mass generation also implies dynamical symmetry breaking in general, as discussed in Ref.[6]. The result has interesting implication to the application to electroweak symmetry breaking, though the focus of the present letter is on the theoretical features.

The details of the calculations involved are very similar to what we have presented in Ref.[6]. In the latter paper, we succeeded in getting the gap equations from first principle supergraph analyses both for the new HSNJL model and the old supersymmetric NJL model[7, 8] with result for the latter case in perfect agreement with the one from the effective field theory analysis [8] and that of the simple NJL limit. We will only sketch the analyses here only highlighting the essential features, for the interest of the general readers. Theorists interested in the details should read carefully also Ref.[6]. Further details together with a fully generic comprehensive analysis of the HSNJL model will be presented in a forthcoming publication [9]. In particular, note that the analyses rely heavily on the formulation of the generating functional, mass parameters, and self-energy amplitudes as superspace quantity as introduced in Ref.[6]. That is only the full superspace analog of usual Minkowski spacetime field theory.

II. DYNAMICAL GENERATION OF MAJORANA MASS

The basic model has two different chiral superfields Φ_+ and Φ_- , presumably carrying different quantum numbers. For instance, they may be different gauge multiplets. The

dimension five four-superfield interaction is given by

$$-\frac{G}{2} \int d^4\theta \Phi_+ \Phi_+ \Phi_- \Phi_- (1 + B\theta^2) \delta^2(\bar{\theta}) . \quad (1)$$

It is really a superpotential term, as indicated by the $\delta^2(\bar{\theta})$, hence holomorphic. In our earlier works [5, 6], the possibility of superfield condensate $\langle \Phi_+ \Phi_- \rangle$ giving rise to Dirac mass term $\mathcal{M} \Phi_+ \Phi_-$ has been investigated. Unlike the dimension six four-superfield interaction, the superpotential term offers also the option of say a condensate of $\langle \Phi_+ \Phi_+ \rangle$ or $\langle \Phi_- \Phi_- \rangle$ giving Majorana mass terms $\mathcal{M}_- \Phi_- \Phi_-$ and $\mathcal{M}_+ \Phi_+ \Phi_+$, respectively. To investigate the option, we proceed similar to the Dirac mass analysis [6], deriving the gap equations for \mathcal{M}_+ and \mathcal{M}_- .

Consider the Lagrangian density written as

$$\begin{aligned} \mathcal{L} = \int d^4\theta \left[\Phi_+^\dagger \Phi_+ (1 - \Delta_+) + \Phi_-^\dagger \Phi_- (1 - \Delta_-) \right. \\ \left. + (\mathcal{M}_+ \Phi_+ \Phi_+ \delta^2(\bar{\theta}) + \mathcal{M}_- \Phi_- \Phi_- \delta^2(\bar{\theta}) + h.c.) \right] + \mathcal{L}_I , \end{aligned} \quad (2)$$

with

$$\mathcal{L}_I = \int d^4\theta \left[-\mathcal{M}_+ \Phi_+ \Phi_+ - \mathcal{M}_- \Phi_- \Phi_- - \frac{G}{2} \Phi_+ \Phi_+ \Phi_- \Phi_- (1 + B\theta^2) \right] \delta^2(\bar{\theta}) + h.c.$$

Here, $\Delta_\pm = \tilde{m}_\pm^2 \theta^2 \bar{\theta}^2$ characterizes soft supersymmetry breaking mass-squared \tilde{m}_\pm^2 for the corresponding scalar field A_\pm and \mathcal{M}_\pm superfield Majorana mass parameter

$$\mathcal{M}_\pm = m_\pm - \theta^2 \eta_\pm , \quad (3)$$

with the supersymmetric Majorana mass m_\pm and its supersymmetry breaking counterpart η_\pm . The gap equations are given by

$$-\mathcal{M}_\pm = \Sigma_{\pm\pm}^{(loop)}(p, \theta^2) \Big|_{\text{on-shell}} , \quad (4)$$

where $\Sigma_{\pm\pm}^{(loop)}$ denotes the lowest order contributions to the proper self-energy from loop diagrams involving the four-superfield interactions. Note that $\Sigma_{++}^{(loop)}$ has contribution involving the Φ_- superfield propagator $\langle T(\Phi_-(1)\Phi_-(2)) \rangle$, and $\Sigma_{--}^{(loop)}$ the propagator $\langle T(\Phi_+(1)\Phi_+(2)) \rangle$. The propagator should include the Majorana masses \mathcal{M}_\pm dependence. The propagators are

given in the same form as the Dirac case of $\langle T(\Phi_+(1)\Phi_-(2)) \rangle$, namely as

$$\begin{aligned} \langle T(\Phi_\pm(1)\Phi_\pm(2)) \rangle &= \frac{i \bar{m}_\pm}{p^2(p^2 + |m_\pm|^2)} \frac{D_1^2}{4} \delta_{12}^4 \\ &\quad - \frac{i}{[(p^2 + |m_\pm|^2 + \tilde{m}_\pm^2)^2 - |\eta_\pm|^2]} \left[\frac{\bar{\eta}_\pm D_1^2 \bar{\theta}_1^2}{4} - \frac{\eta_\pm |m_\pm|^2 D_1^2 \theta_1^2}{4p^2} \right] \delta_{12}^4 \\ &\quad + \frac{i \bar{m}_\pm [\tilde{m}_\pm^2(p^2 + |m_\pm|^2 + \tilde{m}_\pm^2) - |\eta_\pm|^2]}{(p^2 + |m_\pm|^2)[(p^2 + |m_\pm|^2 + \tilde{m}_\pm^2)^2 - |\eta_\pm|^2]} \left[\frac{D_1^2 \theta_1^2 \bar{\theta}_1^2}{4} + \frac{\bar{\theta}_1^2 \theta_1^2 D_1^2}{4} \right] \delta_{12}^4. \quad (5) \end{aligned}$$

We have the gap equations

$$\begin{aligned} m_\pm &= \frac{\bar{\eta}_\mp G}{2} I_2(|m_\mp|^2, \tilde{m}_\mp^2, |\eta_\mp|, \Lambda^2), \\ \eta_\pm &= \bar{m}_\mp G I_1(|m_\mp|^2, \tilde{m}_\mp^2, |\eta_\mp|, \Lambda^2) - \frac{\bar{\eta}_\mp GB}{2} I_2(|m_\mp|^2, \tilde{m}_\mp^2, |\eta_\mp|, \Lambda^2), \quad (6) \end{aligned}$$

where $I_1(|m|^2, \tilde{m}^2, |\eta|, \Lambda^2)$ and $I_2(|m|^2, \tilde{m}^2, |\eta|, \Lambda^2)$ are the same loop integrals as before [6], details of which we will discuss below.

The first thing to note from the gap equation results is that they are almost of exactly the same form as the Dirac case [6]. Actually, if we take identical soft masses $\tilde{m}_\pm^2 = \tilde{m}^2$, we have obviously a symmetric solution relative to Φ_+ and Φ_- , which is exactly the same as the gap equation for the Dirac mass case showed in Ref.[6]. So, nontrivial Majorana masses solution is possible, or as likely as that of the Dirac mass.

III. DIRAC VERSUS MAJORANA MASSES

While the above analysis established the holomorphic four-superfield interaction as being capable of dynamically generating superfield Majorana masses, the result can put the Dirac mass generation scenario under question. After all, true Dirac mass means nonzero $\mathcal{M}\Phi_+\Phi_-$ mass term without \mathcal{M}_+ and \mathcal{M}_- . It is important to note that which mass terms arise have direct implications on what is the symmetry breaking pattern resulted. For example, let us consider Φ_+ and Φ_- each having a Z_2 charge of its own kind while one has the color of quarks and the other of antiquarks. The \mathcal{M}_+ and \mathcal{M}_- mass terms break only the color symmetry, while the Dirac mass term \mathcal{M} breaks both Z_2 symmetries but preserves color. A fully general analysis considering a generic mass matrix for the Φ_+ and Φ_- superfields may have to be performed to answer the important question [9]. However, it is interesting to see that with the gap equation results we have so far one can extract strong indication on

how pure Dirac mass generation may be obtained. The bottom line is, a split in soft masses favors Dirac mass generation over Majorana mass generation. Let us illustrate the story.

In the discussion below, we will compare the gap equation results for the Majorana mass generations discussed above, *i.e.* assuming no off-diagonal $\mathcal{M}\Phi_+\Phi_-$ mass term, and that of the (pure) Dirac mass generation analysis of Ref.[6] where there is also the hidden assumption of no diagonal (Majorana) masses \mathcal{M}_+ and \mathcal{M}_- . This is not the full rigorous way to address the question of what would be the mass matrix of Φ_+ and Φ_- resulted, but can at least give some insight into some qualitative aspect of the question.

As in the above discussions on the Majorana mass case, we focus on the simpler case with $B = 0$. Eliminating G from the gap equations, we get the relations

$$2|m_{\pm}|^2 I_1(|m_{\pm}|^2, \tilde{m}_{\pm}^2, |\eta_{\pm}|, \Lambda^2) = |\eta_{\mp}|^2 I_2(|m_{\mp}|^2, \tilde{m}_{\mp}^2, |\eta_{\mp}|, \Lambda^2) . \quad (7)$$

The integrals are given by ¹

$$\begin{aligned} I_1(|m|^2, \tilde{m}^2, |\eta|, \Lambda^2) &= \frac{1}{16\pi^2} \left[\frac{1}{2}(|m|^2 + \tilde{m}^2) \ln \frac{(\Lambda^2 + |m|^2 + \tilde{m}^2)^2 - |\eta|^2}{(|m|^2 + \tilde{m}^2)^2 - |\eta|^2} \right. \\ &\quad \left. - |m|^2 \ln \frac{(\Lambda^2 + |m|^2)}{|m|^2} + \frac{|\eta|}{2} \ln \frac{(\Lambda^2 + |m|^2 + \tilde{m}^2 + |\eta|)(|m|^2 + \tilde{m}^2 - |\eta|)}{(\Lambda^2 + |m|^2 + \tilde{m}^2 - |\eta|)(|m|^2 + \tilde{m}^2 + |\eta|)} \right] , \\ I_2(|m|^2, \tilde{m}^2, |\eta|, \Lambda^2) &= \frac{1}{32\pi^2} \left[\ln \frac{(\Lambda^2 + |m|^2 + \tilde{m}^2)^2 - |\eta|^2}{(|m|^2 + \tilde{m}^2)^2 - |\eta|^2} \right. \\ &\quad \left. + \frac{|m|^2 + \tilde{m}^2}{|\eta|} \ln \frac{(\Lambda^2 + |m|^2 + \tilde{m}^2 + |\eta|)(|m|^2 + \tilde{m}^2 - |\eta|)}{(\Lambda^2 + |m|^2 + \tilde{m}^2 - |\eta|)(|m|^2 + \tilde{m}^2 + |\eta|)} \right] . \end{aligned} \quad (8)$$

On the $|m| - |\eta|$ plane, the I_2 expression is positive definite while the I_1 expression has maximum value at the origin, given by $\frac{\tilde{m}^2}{16\pi^2} \ln \frac{(\Lambda^2 + \tilde{m}^2)}{\tilde{m}^2}$, which implies that I_1 will be always negative for $\tilde{m}^2 = 0$. What is interesting then is that in the case one of the soft mass vanishes, say $\tilde{m}_+^2 = 0$, $|m_+|$ and $|\eta_-|$ will then be forced to vanish for the above relation to be satisfied. But the tachyonic bound for $|\eta_+|$ [$|\eta_{\pm}| < (|m|^2 + \tilde{m}_{\pm}^2)/2$ for the Majorana case] will then force it to vanish, hence giving also $|m_-| = 0$. That is independent of the value of the coupling.

The argument above shows there will be no nontrivial Majorana masses generated in the case one of the soft masses vanishes. Vanishing, common, soft masses gives no dynamical Dirac mass either [6]. It is easy to see in general, smaller soft supersymmetry breaking mass

¹ The expression is formally equivalent to the earlier form given in Ref.[6], where the \tanh^{-1} function is involved which straightly speaking has a domain of definition problem.

disfavors the dynamical mass generations, requiring a stronger couple G to achieve it, as also explicitly illustrated by the expression for the G threshold given in Ref.[6]. However, when there is a splitting between the soft masses of the two superfields, there is a crucial difference between the Majorana mass and Dirac mass cases. While the supersymmetric part m of each Majorana mass is directly sensitive to the vanishing of the corresponding soft mass, the Dirac mass result is more sensitive to the average of the two soft masses [6]. We have seen that one m vanishes implies all Majorana mass parameters vanishes, at least for $B = 0$. We will show explicitly in the next section that having one vanishing soft mass does not affect the dynamical generation of Dirac mass adversely. Having one soft mass small squeezes the possibility of Majorana mass generation, pushing up the threshold coupling, but has limited effect on the Dirac mass generation.

IV. DIRAC MASS GENERATION WITH ONE OF THE SOFT MASSES VANISHING

Gap equation for the case of Dirac mass generation with two different soft masses has essentially be given in Ref.[6] (see appendix B). We have

$$\begin{aligned} m &= \frac{\bar{\eta}G}{2} I'_2(|m|^2, \tilde{m}_+^2, \tilde{m}_-^2, |\eta|, \Lambda^2) , \\ \eta &= \bar{m}G I'_1(|m|^2, \tilde{m}_+^2, \tilde{m}_-^2, |\eta|, \Lambda^2) - \frac{\bar{\eta}GB}{2} I'_2(|m|^2, \tilde{m}_+^2, \tilde{m}_-^2, |\eta|, \Lambda^2) , \end{aligned} \quad (9)$$

where I'_1 and I'_2 are the loop integrals. The expressions of the integrals as given in Ref.[6], however, have some typos and have not been written in the best form. The integral should be, exactly,

$$\begin{aligned} I'_1(|m|^2, \tilde{m}_+^2, \tilde{m}_-^2, |\eta|, \Lambda^2) &= 2 \int \frac{d^4k}{(2\pi)^4} \frac{\left(\frac{\tilde{m}_+^2 + \tilde{m}_-^2}{2}\right) \left(k^2 + |m|^2 + \frac{\tilde{m}_+^2 + \tilde{m}_-^2}{2}\right) - \left(\frac{\tilde{m}_+^2 - \tilde{m}_-^2}{2}\right)^2 - |\eta|^2}{(k^2 + |m|^2) \left[\left(k^2 + |m|^2 + \frac{\tilde{m}_+^2 + \tilde{m}_-^2}{2}\right)^2 - \left(\frac{\tilde{m}_+^2 - \tilde{m}_-^2}{2}\right)^2 - |\eta|^2 \right]} , \\ I'_2(|m|^2, \tilde{m}_+^2, \tilde{m}_-^2, |\eta|, \Lambda^2) &= \int \frac{d^4k}{(2\pi)^4} \frac{1}{\left[\left(k^2 + |m|^2 + \frac{\tilde{m}_+^2 + \tilde{m}_-^2}{2}\right)^2 - \left(\frac{\tilde{m}_+^2 - \tilde{m}_-^2}{2}\right)^2 - |\eta|^2 \right]} . \end{aligned} \quad (10)$$

What we failed to highlight in that paper are the relations

$$\begin{aligned} I'_1(|m|^2, \tilde{m}_+^2, \tilde{m}_-^2, |\eta|, \Lambda^2) &= I_1(|m|^2, \tilde{m}_{av}^2, |\eta'|, \Lambda^2) , \\ I'_2(|m|^2, \tilde{m}_+^2, \tilde{m}_-^2, |\eta|, \Lambda^2) &= I_2(|m|^2, \tilde{m}_{av}^2, |\eta'|, \Lambda^2) , \end{aligned} \quad (11)$$

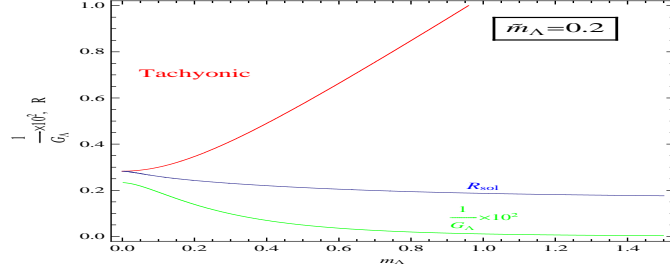


FIG. 1: Illustrative numerical plot of the solution curve on R_Λ - m_Λ plane. R_Λ and m_Λ are corresponding dimensionless parameters for R and $|m|$ normalized to the basic (cut-off) scale Λ . Likewise, $\tilde{m}_\Lambda^2 = \tilde{m}^2/\Lambda^2$. Value of $1/G_\Lambda$ versus m_Λ , $G_\Lambda = |G|\Lambda$ being the normalized coupling, is also given.

where

$$\tilde{m}_{av}^2 = \frac{\tilde{m}_+^2 + \tilde{m}_-^2}{2} \quad (12)$$

is the average value of the soft masses and

$$|\eta'| = \sqrt{|\eta|^2 + \left(\frac{\tilde{m}_+^2 - \tilde{m}_-^2}{2}\right)^2}. \quad (13)$$

The relations are actually easy to appreciate from the physics point of view, by comparing how the various parameters go into the (scalar) mass eigenvalues in the cases with different or identical soft masses.

From the above, one can easily use the properties of the I_1 and I_2 integrals to look at the case to Dirac mass generation with only one of the soft mass vanishes. For example, one can check that since I_1 increases as $|m|$ and $|\eta'|$ decrease, it attains maximum at $|m| = 0$ and $|\eta'|$ minimum of $\left|\frac{\tilde{m}_+^2 - \tilde{m}_-^2}{2}\right|$ which corresponds to $|\eta| = 0$; and the maximum value is given by exactly the same expression as before, namely $\frac{2\tilde{m}_{av}^2}{16\pi^2} \ln \frac{\Lambda^2 + 2\tilde{m}_{av}^2}{2\tilde{m}_{av}^2}$. So, twice the average soft mass here, which equals to the single nonzero soft mass plays the role of the single soft mass relevant for the particular individual Majorana mass equation or the universal soft mass case for Majorana as well as Dirac mass generation analysis. The simple conclusion is having one vanishing soft mass does not kill the Dirac mass generation as it does to the Majorana mass generations.

A further look at the solution curve as given by Eqn.(7) of this Dirac case, $\frac{\bar{\eta}}{2m}I_2 = \frac{\bar{m}}{\eta}I_1$ ($= \frac{1}{G}$), on the plane of $R = |\eta|/|m|$ versus $|m|$ (for fixed Λ and \tilde{m}_{av}^2) is particularly illustrative.

A case example is plotted in Fig. 1. In general, the R decreases monotonically with $|m|$, slowing down to an asymptotical constant value at the $|m| \rightarrow \infty$ limit. On the physics side, it is sensible to restrict all mass parameters not to go beyond the scale of order Λ , which is the model cut-off scale. The asymptotical analysis helps though to show the mathematical features. A careful calculation of the limiting expressions for I_1 and I_2 gives the asymptotic value as $R_\infty = \sqrt{\frac{2\tilde{m}_{av}^2}{3}}$, a constant. For the $|m| \rightarrow 0$ limit, the limiting expressions for I_1 and I_2 gives $I_1 = \tilde{m}_{av}^2 I_2$, hence the constant $R_o = \sqrt{2\tilde{m}_{av}^2}$. The latter is right at the boundary of the inadmissible tachyonic region (for lighter scalar mass eigenstate from A_\pm). The bound itself is given by $R_t = \sqrt{2\tilde{m}_{av}^2 + |m|^2}$, increasing with $|m|$. The solution curve is safely below the tachyonic bound. One can also get the asymptotic expressions for the coupling G , or rather $|G|$. We have $|G_\infty| = 32\sqrt{6}\pi^2|m|^4/\tilde{m}_{av}\Lambda^4$, and the threshold coupling $|G_o| = 16\sqrt{2}\pi^2 \left[\tilde{m}_{av} \ln(1 + \frac{\Lambda}{2\tilde{m}_{av}}) \right]^{-1}$.

V. CONCLUSIONS

We show in this letter that the HSNJL model is capable of dynamically generating Majorana masses for the two superfields involved. This is an important alternative to the Dirac mass generation analyzed earlier. The general question of under what condition a particular mass pattern, or mass matrices, for the two superfields will be resulted becomes an important question to address theoretically. It also has phenomenological implication on whether there is a successful application of the model to electroweak symmetry breaking. We give here a first answer to the complicated question, that a splitting in the soft supersymmetry breaking masses favors Dirac over Majorana mass. In particular, in the limit where one soft mass vanishes, nontrivial Majorana mass is not possible while Dirac mass can still be generated. Hence, one expects a strong coupling within a particular range dictated by the different soft masses values will give a symmetry breaking answer corresponds to the pure Dirac mass.

In the case of the application to the electroweak symmetry breaking [5, 6], the model has a four-superfield interaction involving three different gauge multiplets and hence three soft masses of the top and bottom squark sector. It has already been noted that a split in the soft masses between the two sectors at least will be needed to give the phenomenologically required top and bottom mass ratio. It is encouraging to see that the kind of mass splitting

also disfavors the generation of Majorana masses for the quark superfield multiplets, which of course will break color and electric charge symmetry. We note also that the strong QCD attraction will also play an important role there favoring color singlet vacuum condensates. Full details of all those await further analysis. However, we consider it reasonable to claim, from what we have been able to establish so far, that the HSNJL stands qualitatively viable as a model for electroweak symmetry breaking. We hope to be able to present in a future publication full quantitative results.

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